Light asymmetric dark matter from new strong dynamics

Mads T. Frandsen,* Subir Sarkar, and Kai Schmidt-Hoberg
Rudolf Peierls Centre for Theoretical Physics, University of Oxford, Oxford OX1 3NP, United Kingdom
(Dated: January 15, 2013)

A ~ 5 GeV 'dark baryon' with a cosmic asymmetry similar to that of baryons is a natural candidate for the dark matter. We study the possibility of generating such a state through dynamical electroweak symmetry breaking, and show that it can share the relic baryon asymmetry via sphaleron interactions, even though it has no electroweak interactions. The scattering cross-section on nucleons, estimated in analogy to QCD, is within reach of underground direct detection experiments.

PACS numbers: 95.35.+d,12.60.Nz

ASYMMETRIC DARK MATTER

A particle-antiparticle asymmetry in dark matter, similar to that in baryons, would provide a natural link between their observed abundances. A classic example of such *asymmetric* dark matter (ADM) is the lightest neutral technibaryon (TB) [1–3] with a mass of $\mathcal{O}(1)$ TeV in technicolour (TC) models of electroweak (EW) symmetry breaking [4]. Other techni-interacting massive particles (TIMPs) [5, 6] have been considered, some of which are pseudo Nambu-Goldstone bosons (pNGBs) of the TC interactions, hence lighter with mass of $\mathcal{O}(100)$ GeV.

While gravitational instability in collisionless cold dark matter (CDM) explains structure formation on large scales well, the predicted substructure on galactic scales is at variance with observations suggesting that CDM may be *self-interacting* [7]. If ADM arises in a strongly coupled theory as a *composite* particle χ then it would naturally have self-interactions, with a cross-section large enough to address the small scale structure issue if their strong-interaction scale is of $\mathcal{O}(\mathrm{GeV})$ (assuming the scaling resembles that in QCD). The self-annihilation cross-section would naturally be of the same order, ensuring that no significant symmetric abundance survives from the early universe. The relic abundance is then given simply by $\Omega_\chi = (m_\chi \mathcal{N}_\chi / m_\mathrm{B} \mathcal{N}_\mathrm{B}) \Omega_B$ where $\mathcal{N}_{\mathrm{B},\chi}$ are the respective asymmetries of baryons and ADM (*e.g.* $\mathcal{N}_\mathrm{B} \equiv (n_\mathrm{B} - n_{\bar{\mathrm{B}}})/(n_\mathrm{B} + n_{\bar{\mathrm{B}}})$). Now if some process ensures $\mathcal{N}_\chi \sim \mathcal{N}_\mathrm{B}$ then the observed cosmological dark matter abundance is realised for a ~ 5 GeV χ particle. Excitingly, recent signals in the underground direct detection experiments DAMA [8] and CoGeNT [9] have been interpreted in terms of such light dark matter [10] and generated renewed interest in GeV scale ADM from new strong dynamics [11]. While subsequent experiments like XENON100 [12] and CDMSII [13] have not confirmed these claims, their results still allow ~ 5 GeV CDM with a spin-independent scattering cross-section on nucleons as high as $\sim 10^{-39}$ cm^2. Moreover such particles can be accreted by the Sun in large numbers (since they cannot annihilate) and affect heat transport so as to measurably alter the fluxes of low energy neutrinos [14].

In this letter we consider a mechanism for generating light asymmetric dark matter by extending the standard $SU(3)_\mathrm{c} \otimes SU(2)_\mathrm{L} \otimes U(1)_Y$ model (SM) to include a new strong interaction exhibiting two sectors S_1, S_2:
i) S_1 breaks EW symmetry dynamically at a scale Λ_1 and the composite spectrum includes $\mathcal{O}(1)$ TeV mass particles TB carrying a new global $U(1)_\mathrm{TB}$, just like 'technibaryons' in technicolour [4]. The constituents of TB carry weak quantum numbers such that TB couples to the fermion number violating EW sphaleron interactions which distribute any pre-existing fermion asymmetry between baryons and TB's, ensuring that:

$$\mathcal{N}_\mathrm{TB} \sim \mathcal{N}_\mathrm{B}. \tag{1}$$

ii) S_2 is uncharged under the SM and becomes strongly interacting at a scale Λ_2 of $\mathcal{O}(\mathrm{GeV})$; its spectrum includes (composite) few GeV mass particles χ also carrying the global $U(1)_\mathrm{TB}$ quantum number. The two sectors are coupled via operators that induce $U(1)_\mathrm{TB}$ preserving fast decays: TB $\to \chi + X$. This interaction keeps χ in equilibrium with T down to $T \lesssim 0.1 m_\mathrm{TB}$, *i.e.* below the temperature $T_\mathrm{sph} \sim m_W$ at which the sphaleron interactions 'freeze-out'. Thus the sphaleron induced asymmetry in TB is converted into a similar asymmetry in χ:

$$\mathcal{N}_\chi \sim \mathcal{N}_\mathrm{TB} \ . \tag{2}$$

This naturally connects the relic density of light ADM to the relic density of baryons via sphalerons, *without* requiring χ (or its constituents) to carry EW quantum numbers. Thus experimental constraints on such light particles from the W decay width *etc* are not relevant.

In fact all that is necessary to generate light ADM via sphalerons is EW charged states at the weak scale which can decay rapidly into ADM. However we consider it more appealing to link ADM to dynamical EW symmetry breaking and consider the possibility that the two sectors above arise from a *single* strongly interacting (technicolour) extension of the SM which develops two different dynamical scales $\Lambda_{1,2}$.

This can happen *e.g.* if the theory contains fermions in *different* representations of the TC gauge group since the critical value α_c of the coupling which breaks chiral symmetry depends on the quadratic Casimir C_2 of the

representation (using a simple one-gluon exchange estimate) [15, 16]. This possibility has been considered in 'two-scale' TC [17] — a variant of the earlier idea [16] that if QCD contains fermions in a higher-dimensional representation, then these might dynamically break the EW symmetry at the correct scale. A second possibility is if large four-fermion operators are present as in the gauged Nambu-Jona-Lasanio (NJL) model [18]. If the four-fermion coupling is sufficiently large, this interaction can drive chiral symmetry breaking, allowing a much smaller value of the critical gauge coupling as in 'top quark condensation' [19] or 'topcolour' [20] models. If only some of the fermions participate in the four-fermion interactions, a scale separation can arise. SM singlet techni-fermions have been introduced earlier [21] in models of 'Minimal Walking Technicolour' [23] and 'Conformal Technicolor' [24], in order to achieve (near-)conformal or 'walking' dynamics, while still maintaining a minimal sector which breaks EW symmetry.

TWO SCALES FROM A STRONGLY INTERACTING THEORY

We consider a non-Abelian gauge theory with gauge group G: N_1 fermions transforming according to a representation $\mathcal{R}_1$ of G are gauged under the EW symmetry, and N_2 fermions transforming according to a representation $\mathcal{R}_2$ of G are SM singlets.

Using the ladder approximation to the Schwinger-Dyson equations, the critical value of the coupling for chiral symmetry breaking is [25]:

$$\alpha_{\rm c} = \frac{\pi}{3C_2(\mathcal{R})}. \tag{3}$$

We now take the gauge group to be $SU(N_{\rm C})$ and consider representations such that $C_2(\mathcal{R}_1) \geq C_2(\mathcal{R}_2)$. Integrating the one-loop beta-function $\beta(\alpha) = -(\frac{\alpha^2}{2\pi}\beta_0 + \frac{\alpha^3}{8\pi^2}\beta_1 + \ldots)$ from Λ_1 to Λ_2 yields the ratio of the scales:

$$\frac{\Lambda_1}{\Lambda_2} \simeq \exp\left[\frac{2\pi}{\beta_0(\mathcal{R}_2)}\left(\alpha_{\rm c}(\mathcal{R}_1)^{-1} - \alpha_{\rm c}(\mathcal{R}_2)^{-1}\right)\right] , \tag{4}$$

Since $\Lambda_1 \geq \Lambda_2$,or equivalently $\alpha_{\rm c}(\mathcal{R}_1) \leq \alpha_{\rm c}(\mathcal{R}_2)$, the fermions in the representation $\mathcal{R}_1$ are, in this approximation, decoupled below Λ_1, so only $\beta_0(\mathcal{R}_2)$ appears in the exponent. If $\beta_0(\mathcal{R}_2)$ and $\alpha_{\rm c}$ are small then the scale separation can be large, *i.e.* we can have $\Lambda_1 \sim \Lambda_{\rm TC}$ and $\Lambda_2 \sim \Lambda_{\rm ADM}$. Our estimate of $\alpha_{\rm c}$ should be compared to the two-loop fixed point value of the coupling:

$$\alpha_* = -4\pi\frac{\beta_0}{\beta_1}. \tag{5}$$

If $\alpha_* < \alpha_{\rm c}$ the theory will run to an infra-red fixed point before triggering chiral symmetry breaking. The lower boundary of the conformal window is thus identified by demanding $\alpha_*(\mathcal{R}_1, \mathcal{R}_2) = \alpha_{\rm c}(\mathcal{R}_1)$.

For the fermions transforming under $\mathcal{R}_1$ to break the EW symmetry at Λ_1 they must be charged under the EW gauge group. The minimal choice (dictated also by constraints from EW precision measurements [6]) is $N_1 = 2$ Dirac flavors with the left-handed Weyl spinors arranged in an $SU(2)_{\rm L} \otimes U(1)_Y$ weak doublet $Q_{\rm L}$, along with N_2 SM singlet Dirac flavors λ:

$$\begin{aligned} Q^a_{\rm L} &= \begin{pmatrix} U^a \\ D^a \end{pmatrix}_{\rm L}, \; U^a_{\rm R}, \; D^a_{\rm R}, \; \lambda^b \\ a &= 1, \ldots d(\mathcal{R}_1), \; b = 1, \ldots d(\mathcal{R}_2). \end{aligned} \tag{6}$$

The condensate $\langle U_{\rm L} U^*_{\rm R} + D_{\rm L} D^*_{\rm R} + {\rm h.c.}\rangle$ breaks EW symmetry and the TBs are made out of the Us and Ds while the dark matter candidate χ is made of the λs (and is a fermion or boson depending on $N_{\rm C}$). The unbroken global symmetries of the TC interactions (depending on $\mathcal{R}_{1,2}$) keep χ stable and are discussed further below.

In Fig. 1 we show the conformal phase-diagram in the $(N_2, N_{\rm C})$ plane, as well as the corresponding scale separation, with N_2 taking the value along the lower boundary of the conformal window. It is seen that we *cannot* achieve large scale separations without having to increase N_2 (*i.e.* reduce $\beta_0(\mathcal{R}_2)$) to values where the theory is actually IR-conformal. Below Λ_1, where the $\mathcal{R}_1$ fermions decouple, the coupling simply runs too fast.

We consider now the effects of four-fermion operators, expected in any case in a more complete theory *e.g.* when embedding the technicolour gauge group in an extended technicolour (ETC) gauge group [26], which is required to communicate EW breaking to the SM fermions without introducing fundamental scalars. It is assumed to break at some high scale $\Lambda_{\rm ETC}$, so in integrating out the heavy ETC gauge bosons, four-fermion vector currents appear at the TC scale. This can also alleviate the tension with the S-parameter [27]. Four-fermion operators can also be induced by the non-perturbative dynamics of walking technicolour itself [28].

We adopt the gauged NJL model [18] as a representative simple theory with four-fermion interactions which leads to chiral symmetry breaking. While the four-fermion operator coupling is of the form $VV - AA$ here (where V and A denote vector and axial vector currents), the relation between chiral symmetry breaking on the one hand and the gauge and four-fermion coupling on the other hand is qualitatively the same in other cases as well. The gauged NJL model for the fermions Q transforming under $\mathcal{R}_1$ is defined by

$$\begin{aligned} \mathcal{L} &= \bar{Q} i \not{D} Q - \frac{1}{4}{\rm Tr} F^a_{\mu\nu} F^{a\mu\nu} \\ &+ \frac{4\pi^2 g_1}{\Lambda^2 N_1 d(\mathcal{R}_1)} \left[(\bar{Q}Q)^2 + (\bar{Q} i\gamma_5 T^a Q)^2\right], \end{aligned} \tag{7}$$

with g_1 the dimensionless four-fermion coupling and Λ the effective cut-off of the the NJL model. The critical

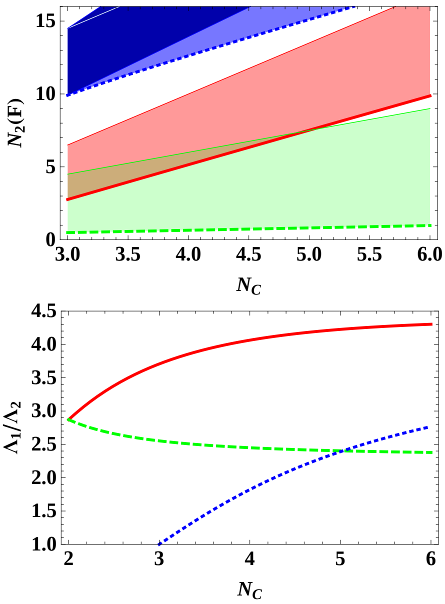

FIG. 1: Upper panel: The conformal window for $SU(N_C)$ gauge theories with $N_2(F)$ fermions in the fundamental representation ($\mathcal{R}_2 = F$), and 2 flavours in the (top to bottom): fundamental (dark blue), two-index antisymmetric (light blue), two-index symmetric (red), and adjoint (green) representations. Lower panel: The corresponding scale separation using the value of $N_2(F)$ at the lower boundary of the corresponding conformal window for the: two-index antisymmetric (blue dotted), adjoint (green dashed), and two-index symmetric (red solid) representations.

line in the (α, g_1)-plane is now given by [29, 30]:

$$\alpha_c(\mathcal{R}_1, g_1) = \begin{cases} 4\left(\sqrt{g_1} - g_1\right) \times \alpha_c(\mathcal{R}_1) & \text{for } \frac{1}{4} < g_1 \leq 1\,, \\ \alpha_c(\mathcal{R}_1) & \text{for } 0 \leq g_1 < \frac{1}{4}\,. \end{cases}$$

The critical gauge coupling is thus reduced and in the limit $g_1 \to 1$, chiral symmetry breaking is driven purely by the four-fermion interaction. The lower boundary of the conformal window, found by imposing $\alpha_*(\mathcal{R}_1, \mathcal{R}_2) = \alpha_c(\mathcal{R}_1, g_1)$, now changes [31] as shown in Fig. 2 for the case of $SU(3)$ gauge theories with $N_1 = 2$ and $\mathcal{R}_1$ in different representations. It is seen how the conformal window shrinks as $g_1 \to 1$, allowing the addition of more matter transforming under $\mathcal{R}_2$ while still having a theory that breaks chiral symmetry. Fig. 2 also shows that the scale separation Λ_2/Λ_1 can now be substantial.

Although we have used rather simple approximations, this shows qualitatively how a significant scale separation can be achieved within a single strongly interacting theory. Other dynamical frameworks may allow large scale separations *e.g.* models with fundamental scalars which condense at Λ_1 and thereby have a similar effect to four-fermion operators. Another possibility is chiral gauge theories where some fermions and some gauge bosons condense at Λ_1 resulting in a slowly evolving coupling below Λ_1 and a large scale separation.

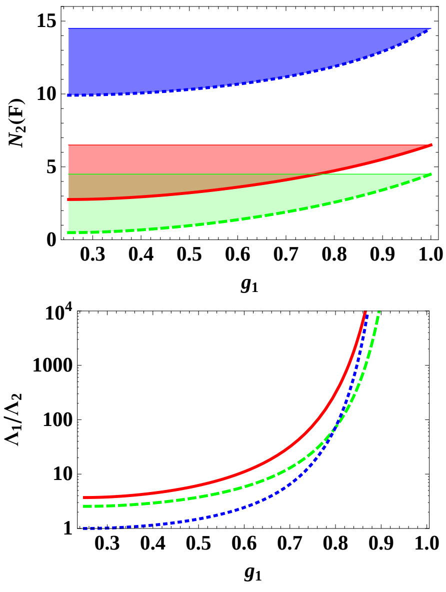

FIG. 2: As in Fig. 1, but now as a function of g_1 (7). Here the fundamental and antisymmetric representations are identical.

PHENOMENOLOGY

The phenomenology will depend on the specific TC gauge group and the fermion representations breaking EW symmetry, but there are some generic implications.

If $\mathcal{R}_1 = \mathcal{R}_2$ (both complex), there is a single global $U(1)_{\rm TB}$ keeping χ stable. We then expect to have a sizeable Yukawa-type interaction between the heavy technibaryon TB, the dark matter particle χ and *e.g.* the composite Higgs which would keep the TB and χ particles in thermal equilibrium in the early universe. However if $\mathcal{R}_1 \neq \mathcal{R}_2$, additional global $U(1)$'s in the theory can prevent such operators from arising; in that case we would require additional interactions such as ETC breaking these $U(1)$'s such that only the technibaryon in the low scale sector remains stable [2]; an interesting example of interactions allowing this has been given [32]. We discuss explicit models elsewhere but comment here on the implications for dark matter detection experiments.

The symmetric component of χ can effectively annihilate into states which do not carry the $U(1)_{\rm TB}$ (*e.g.* technipions) and can subsequently decay to SM states. The asymmetric component of χ cannot annihilate so there are no indirect signatures (since decays are highly suppressed compared to TeV scale ADM [33]).

We expect isospin-0 scalar and vector mesons like a 'techni-sigma' (or composite Higgs) and a 'techni-omega' in both sectors. The 'techni-omega' will mix with the SM hypercharge field after EW symmetry breaking and induce couplings of χ to the SM sector, just as we expect the 'techni-sigma' to couple to SM fermions via effective Yukawa couplings (induced *e.g.* by ETC). This can lead to exciting signals at the LHC [34].

The spin-idependent elastic scattering cross-section on nucleons from either scalar or vector boson exchange is

$$\sigma \sim g_\chi^2 g_q^2 \frac{\mu^2}{m^4}, \tag{8}$$

which is $\sim (10^{-32} - 10^{-30}) g_q^2$ cm^2 for a mediator mass of $m \sim 5 - 15$ GeV, a reduced mass $\mu \sim 1$ GeV and a coupling $g_\chi \sim 1$ between χ and the mediator. The coupling g_q arises from mixing between the light and heavy states, thus parametrically $g_q \sim \Lambda_2/\Lambda_1$, and in addition it is proportional to small couplings — the $U(1)_Y$ coupling g' and fermion hypercharges in the 'techni-omega' case, and the light SM fermion Yukawa couplings in the composite Higgs case. Hence we expect $g_q \lesssim 10^{-4}$, so the cross-section is within reach of direct detection experiments which are sensitive to nuclear recoil energies of $\mathcal{O}$(keV) characteristic of $\sim$5 GeV ADM.

While pNGB's themselves may carry $U(1)_{\rm TB}$ [5] and provide another way of generating light ADM in simple models of dynamical EW symmetry breaking, their self-interactions are derivatively suppressed at low energies. By contrast, the 'dark baryon' ADM state χ considered here is expected to have large self-interactions and thus interesting astrophysical signatures [7, 14].

Acknowledgements

MTF acknowledges a VKR Foundation Fellowship. We thank John March-Russell and, especially, Francesco Sannino for stimulating discussions, and the European Research and Training Network "Unification in the LHC era" (PITN-GA-2009-237920) for partial support.

---

* Electronic address: m.frandsen1@physics.ox.ac.uk

[1] S. Nussinov, Phys. Lett. B **165**, 55 (1985).

[2] R.S. Chivukula & T.P. Walker, Nucl. Phys. B **329**, 445 (1990).

[3] S.M. Barr, R.S. Chivukula & E. Farhi, Phys. Lett. B **241**, 387 (1990).

[4] S. Weinberg, Phys. Rev. D **19**, 1277 (1979); L. Susskind, Phys. Rev. D **20**, 2619 (1979).

[5] S.B. Gudnason, C. Kouvaris & F. Sannino, Phys. Rev. D **73**, 115003 (2006); C. Kouvaris, Phys. Rev. D **76**, 015011 (2007). F. Sannino & R. Zwicky, Phys. Rev. D **79**, 015016 (2009); R. Foadi, M.T. Frandsen & F. Sannino, Phys. Rev. D **80**, 037702 (2009); M.T. Frandsen & F. Sannino, Phys. Rev. D **81**, 097704 (2010); A. Belyaev *et al.*, Phys. Rev. D **83** (2011) 015007.

[6] T.A. Ryttov & F. Sannino, Phys. Rev. D **78**, 115010 (2008).

[7] D.N. Spergel & P.J. Steinhardt, Phys. Rev. Lett. **84**, 3760 (2000); C.J. Hogan & J.J. Dalcanton, Phys. Rev. **D62**, 063511 (2000); R. Dave *et al.*, Astrophys. J. **547**, 574 (2001).

[8] R. Bernabei *et al.*, Eur. Phys. J. C **67**, 39 (2010)

[9] C.E. Aalseth *et al.* arXiv:1002.4703.

[10] J. Kopp, T. Schwetz & J. Zupan, JCAP **1002**, 014 (2010); M. Farina, D. Pappadopulo & A. Strumia, Phys. Lett. B **688**, 329 (2010); A.L. Fitzpatrick, D. Hooper & K.M. Zurek, Phys. Rev. D **81**, 115005 (2010); C. Savage, *et al*, Phys. Rev. D **83**, 055002 (2011).

[11] D.B. Kaplan, Phys. Rev. Lett. **68**, 741 (1992); D. Hooper, J. March-Russell & S.M. West, Phys. Lett. B **605**, 228 (2005); D.E. Kaplan, M.A. Luty & K.M. Zurek, Phys. Rev. D **79**, 115016 (2009); H. An *et al.*, JHEP **1003**, 124 (2010); T. Cohen *et al.*, Phys. Rev. D **82**, 056001 (2010); J. Shelton & K.M. Zurek, Phys. Rev. D **82**, 123512 (2010) N. Haba & S. Matsumoto, arXiv:1008.2487; M.R. Buckley & L. Randall, arXiv:1009.0270; P.H. Gu *et al*, arXiv:1009.2690; M. Blennow *et al*, JHEP **1103**, 014 (2011); L.J. Hall, J. March-Russell & S.M. West, arXiv:1010.0245; B. Dutta & J. Kumar, arXiv:1012.134; A. Falkowski, J.T. Ruderman & T. Volansky, arXiv:1101.4936; J.J. Heckman & S.J. Rey, arXiv:1102.5346; M.L. Graesser, I.M. Shoemaker, L. Vecchi, arXiv:1103.2771.

[12] E. Aprile *et al.* Phys. Rev. Lett. **105**, 131302 (2010).

[13] D.S. Akerib *et al.* Phys. Rev. **D82**, 122004 (2010).

[14] M.T. Frandsen & S. Sarkar, Phys. Rev. Lett. **105**, 011301 (2010); D.T. Cumberbatch *et al.*,Phys. Rev. D **82**, 103503 (2010); M. Taoso *et al.*, Phys. Rev. D **82**, 083509 (2010).

[15] S. Raby, S. Dimopoulos & L. Susskind, Nucl. Phys. B **169**, 373 (1980).

[16] W.J. Marciano, Phys. Rev. D **21**, 2425 (1980).

[17] K.D. Lane & E. Eichten, Phys. Lett. B **222**, 274 (1989).

[18] Y. Nambu & G. Jona-Lasinio, Phys. Rev. **122**, 345 (1961); K. Yamawaki, arXiv:hep-ph/9603293.

[19] V.A. Miransky, M. Tanabashi & K. Yamawaki, Phys. Lett. B **221**, 177 (1989).

[20] C.T. Hill, Phys. Lett. B **345**, 483 (1995)

[21] D.D. Dietrich, F. Sannino & K. Tuominen, Phys. Rev. D **72**, 055001 (2005);

[22] M.A. Luty, JHEP **0904**, 050 (2009).

[23] F. Sannino & K. Tuominen, Phys. Rev. D **71**, 051901 (2005).

[24] M.A. Luty & T. Okui, JHEP **0609**, 070 (2006)

[25] T. Appelquist, K.D. Lane & U. Mahanta, Phys. Rev. Lett. **61**, 1553 (1988).

[26] S. Dimopoulos & L. Susskind, Nucl. Phys. B **155**, 237 (1979); E. Eichten & K.D. Lane, Phys. Lett. B **90**, 125 (1980).

[27] R.S. Chivukula *et al.*, Phys. Rev. D **71**, 115001 (2005)

[28] M. Kurachi, R. Shrock & K. Yamawaki, Phys. Rev. D **76**, 035003 (2007)

[29] K.i. Kondo, H. Mino & K. Yamawaki, Phys. Rev. D **39**, 2430 (1989).

[30] T. Appelquist *et al*, Proc. 12th Johns Hopkins Workshop on Current Problems in Particle Theory (1988) p.197.

[31] H.S. Fukano & F. Sannino, Phys. Rev. D **82**, 035021 (2010)

[32] F. Wilczek & A. Zee, Phys. Rev. D **16**, 860 (1977).

[33] E. Nardi, F. Sannino & A. Strumia, JCAP **0901**, 043 (2009)

[34] R.S. Chivukula & M. Golden, Phys. Rev. D **41**, 2795 (1990); A.R. Zerwekh, Eur. Phys. J. C **46**, 791 (2006); R. Barbieri, S. Rychkov & R. Torre, Phys. Lett. B **688**, 212 (2010).